Sneutrino Hybrid Inflation in Supergravity

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We propose a hybrid inflation scenario in which the singlet sneutrino, the superpartner of the right-handed neutrino, plays the role of the inflaton. We study a minimal model of sneutrino hybrid inflation in supergravity, where we find a spectral index $n_\mathrm{s} \approx 1 + 2\gamma$ with $|\gamma| \lesssim 0.02$, and predict a running spectral index $|\mathrm{d}n_\mathrm{s}/\mathrm{d}\ln k| \ll |\gamma|$ and a tensor-to-scalar ratio $r \ll \gamma^2$ for field values well below the Planck scale. In our scenario, the baryon asymmetry of our universe can be explained via non-thermal leptogenesis and a low reheat temperature $T_\mathrm{RH} \approx 10^6$ GeV can be realized.

Introduction — Although inflation has become a widely accepted paradigm for solving the flatness and horizon problems of the early universe, the relation of inflation to particle physics remains unclear. Even though there are many models of inflation, there are not so many particle physics candidates for the scalar field responsible for inflation, the so called inflaton field [1]. The main reason for this is that the inflaton is required to satisfy slow-roll conditions which are in general difficult to reconcile with the known couplings of particle physics candidates. The experimental discovery of neutrino mass and mixing, when combined with the ideas of the see-saw mechanism [2] and supersymmetry [3], gives a new perspective on this problem. In order to generate the observed neutrino masses within a see-saw extended version of the minimal supersymmetric standard model (MSSM), right-handed neutrinos are typically introduced and small neutrino masses arise naturally from the see-saw mechanism. Among the particles of this extended MSSM, the singlet sneutrinos, the superpartners of the right-handed neutrinos, become attractive candidates, in fact the only candidates, for playing the role of the inflaton. Motivated by such considerations, the possibility of chaotic (large field) inflation with a sneutrino inflaton [4] has recently been revisited [5]. An alternative to chaotic inflation is hybrid inflation [6], which, in contrast to chaotic inflation, involves field values well below the Planck scale. Hybrid inflation is thereby promising for connecting inflation to particle physics [7].

In this Letter we suggest that one (or more) of the singlet sneutrinos $\tilde{N}_i$, where $i = 1, 2, 3$ is a family index, could be the inflaton of hybrid inflation. We present a minimal model of sneutrino hybrid inflation in supergravity, and find the spectral index $n_\mathrm{s} \approx 1 + 2\gamma$ with $|\gamma| \lesssim 0.02$, $|\mathrm{d}n_\mathrm{s}/\mathrm{d}\ln k| \ll |\gamma|$ and the tensor-to-scalar ratio $r \ll \gamma^2$. We shall furthermore show how the baryon asymmetry of our universe can be explained via non-thermal leptogenesis [8, 9]. A low reheat temperature $T_\mathrm{RH} \approx 10^6$ GeV, consistent with the gravitino constraints

in some supergravity theories [10], can be realized with values of the neutrino Yukawa couplings consistent with first family quark and lepton Yukawa couplings in Grand Unified Theories (GUTs).

The Model — Consider the superpotential [11]

$$\mathcal{W} = \kappa \hat{S} \left(\frac{\hat{\phi}^4}{M'^2} - M^2 \right) + \frac{(\lambda_N)_{ij}}{M_*} \hat{N}_i \hat{N}_j \, \hat{\phi}\hat{\phi} + \dots, \quad (1)$$

where κ and $(\lambda_N)_{ij}$ are dimensionless Yukawa couplings and M, M' and M_* are three independent mass scales. The superfields $\hat{N}_i$, $\hat{\phi}$ and $\hat{S}$ contain the following bosonic components, respectively: the singlet sneutrino inflaton $\tilde{N}$ [14], which is non-zero during inflation; the so-called waterfall field ϕ, which is held at zero during inflation but which develops a non-zero vacuum expectation value (vev) after inflation; and the singlet field S which is held at zero during and after inflation [15]. The form of $\mathcal{W}$ in Eq. (1) can be understood as follows. The first term on the right-hand side serves to fix the vev of the waterfall field after inflation and contributes a large vacuum energy to the potential during inflation. We have chosen the waterfall superfield to appear in this term as $\hat{\phi}^4/M'^2$ instead of $\hat{\phi}^2$ in order to allow a Z_4 discrete symmetry to prevent explicit singlet (s)neutrino masses [17]. $\mathcal{W}$ is also compatible with a $U(1)_\mathrm{R}$-symmetry under which $\mathcal{W}$ and $\hat{S}$ each carry unit R-charge, while the charge of $\hat{N}$ is $1/2$. Under suitable conditions the discrete subgroup of this symmetry acts as matter parity [16]. The second term on the right-hand side of Eq. (1) allows the sneutrino inflaton to give a positive mass squared for the waterfall field during inflation, which fixes its vev at zero as long as $|\tilde{N}|$ is above a critical value. After inflation, when the waterfall field acquires its non-zero vev, the same term yields the masses of the singlet (s)neutrinos.

With non-zero F-terms during inflation, the Kähler potential can contribute significantly to the scalar potential. Since the field values of the inflaton are well below the reduced Planck scale $m_\mathrm{P} = 1/\sqrt{8\pi G_\mathrm{N}}$, we can consider

an expansion in powers of $1/m_\mathrm{P}^2$ [18]:

$$\mathcal{K} = |\hat{S}|^2 + |\hat{\phi}|^2 + |\hat{N}|^2 + \kappa_S \frac{|\hat{S}|^4}{4m_\mathrm{P}^2} + \kappa_N \frac{|\hat{N}|^4}{4m_\mathrm{P}^2} + \kappa_\phi \frac{|\hat{\phi}|^4}{4m_\mathrm{P}^2}$$

$$+ \kappa_{S\phi} \frac{|\hat{S}|^2|\hat{\phi}|^2}{m_\mathrm{P}^2} + \kappa_{SN} \frac{|\hat{S}|^2|\hat{N}|^2}{m_\mathrm{P}^2} + \kappa_{N\phi} \frac{|\hat{N}|^2|\hat{\phi}|^2}{m_\mathrm{P}^2} + \ldots \ (2)$$

where the dots indicate higher order terms and additional terms for the other fields in the theory. With a non-canonical Kähler potential as above, the field S acquires a large mass which holds it at zero during inflation. We will neglect radiative corrections to the potential in the following, which are generically subdominant in our model.

The Potential — We now analyze the scalar potential for the model defined by the superpotential $\mathcal{W}$ of Eq. (1) and the Kähler potential $\mathcal{K}$ of Eq. (2). The F-term contributions to the scalar potential are given by:

$$V_\mathrm{F} = e^{\mathcal{K}/m_\mathrm{P}^2} \left[K_{ij}^{-1} D_{z_i}\mathcal{W} D_{z_j^*}\mathcal{W}^* - 3m_\mathrm{P}^{-2}|\mathcal{W}|^2 \right], \ (3)$$

with z_i being the bosonic components of the superfields $\hat{z}_i \in \{\hat{N}, \hat{\phi}, \hat{S}, \ldots\}$ [20] and where we have defined

$$D_{z_i}\mathcal{W} := \frac{\partial\mathcal{W}}{\partial z_i} + m_\mathrm{P}^{-2}\frac{\partial\mathcal{K}}{\partial z_i}\mathcal{W}, \ K_{ij} := \frac{\partial^2\mathcal{K}}{\partial z_i \partial z_j^*} \quad (4)$$

and $D_{z_j^*}\mathcal{W}^* := (D_{z_j}\mathcal{W})^*$. Since we assume that $\tilde{N}, \phi$ and S are effective gauge singlets at the energy scales under consideration, there are no relevant D-term contributions. From Eqs. (1) and (2), with canonically normalized fields, we obtain the potential

$$V = \kappa^2 \left(\frac{|\phi|^4}{M'^2} - M^2 \right)^2 \left(1 + (1 - \kappa_{S\phi})\frac{|\phi|^2}{m_\mathrm{P}^2} + \right.$$

$$+ \ (1 - \kappa_{SN})\frac{|\tilde{N}|^2}{m_\mathrm{P}^2} - \kappa_S \frac{|S|^2}{m_\mathrm{P}^2} \Bigg)$$

$$+ \frac{4\lambda_N^2}{M_*^2}\left(|\tilde{N}|^4|\phi|^2 + |\tilde{N}|^2|\phi|^4 \right) + \ldots , \quad (5)$$

where we have shown only the leading order terms and the terms essential for our analysis.

Sneutrino Hybrid Inflation — Writing the potential of Eq. (5) in terms of real fields $\tilde{N}_\mathrm{R} = \sqrt{2}|\tilde{N}|$, $\phi_\mathrm{R} = \sqrt{2}|\phi|$ and $S_\mathrm{R} = \sqrt{2}|S|$, we obtain

$$V = \kappa^2 \left(\frac{\phi_\mathrm{R}^4}{4M'^2} - M^2 \right)^2 \left(1 - \beta\frac{\phi_\mathrm{R}^2}{2m_\mathrm{P}^2} + \gamma\frac{\tilde{N}_\mathrm{R}^2}{2m_\mathrm{P}^2} \right.$$

$$\left. - \kappa_S \frac{S_\mathrm{R}^2}{2m_\mathrm{P}^2} \right) + \frac{\lambda_N^2}{2M_*^2}\left(\tilde{N}_\mathrm{R}^4\phi_\mathrm{R}^2 + \tilde{N}_\mathrm{R}^2\phi_\mathrm{R}^4 \right) + \ldots , \ (6)$$

where we have defined

$$\beta := \kappa_{S\phi} - 1 \quad (> 0 \text{ for inflation to end}), \quad (7)$$

$$\gamma := 1 - \kappa_{SN}. \quad (8)$$

During inflation, the waterfall field ϕ_R has a zero vev and the potential is dominated by the vacuum energy $V_0 = \kappa^2 M^4$. This false vacuum during inflation is stable as long as the mass squared for the waterfall field ϕ_R is positive. From Eq. (6) we obtain the requirement

$$m_{\phi_\mathrm{R}}^2 = \lambda_N^2 \frac{\tilde{N}_\mathrm{R}^4}{M_*^2} - \beta\frac{\kappa^2 M^4}{m_\mathrm{P}^2} > 0. \quad (9)$$

Inflation ends when $m_{\phi_\mathrm{R}}^2$ becomes negative, i.e. ϕ_R develops a tachyonic instability and rolls rapidly to its global minimum at $\langle\phi_\mathrm{R}\rangle = \sqrt{2M'M}$. Clearly, this requires $\beta > 0$, as already indicated in Eq. (7). More precisely, inflation ends by a second order phase transition when the field value of the inflaton drops below the critical value $\tilde{N}_{\mathrm{R}c}$ given by

$$\tilde{N}_{\mathrm{R}c}^2 = \sqrt{\beta}\frac{\kappa}{\lambda_N}\frac{M^2 M_*}{m_\mathrm{P}}. \quad (10)$$

From Eq. (6) we see that S_R can be set to zero during inflation if we take e.g. $\kappa_S < -1/3$, such that S_R gets a mass term larger than the Hubble parameter $H \approx \sqrt{V_0}/(\sqrt{3}m_\mathrm{P})$. With $\phi_\mathrm{R} = S_\mathrm{R} = 0$, the part of the scalar potential relevant for the evolution of the singlet sneutrino inflaton $\tilde{N}_\mathrm{R}$ during inflation is given by

$$V = \kappa^2 M^4 \left(1 + \gamma\frac{\tilde{N}_\mathrm{R}^2}{2m_\mathrm{P}^2} + \delta\frac{\tilde{N}_\mathrm{R}^4}{4m_\mathrm{P}^4} \right) + \ldots , \quad (11)$$

where we have included the next-to-leading order term proportional to $\delta = \frac{1}{2} + \kappa_{SN}^2 - \kappa_{SN}\kappa_N + \frac{5}{4}\kappa_N + \ldots$. The parameter γ in the scalar potential controls the mass of the inflaton. Furthermore, compared to the term proportional to γ, the term proportional to δ is suppressed by $\tilde{N}_\mathrm{R}^2/m_\mathrm{P}^2$ and will be neglected. The slow-roll parameters are given by

$$\epsilon := \frac{m_\mathrm{P}^2}{2}\left(\frac{V'}{V}\right)^2 \simeq \frac{(\delta\tilde{N}_\mathrm{R}^3 + m_\mathrm{P}^2\gamma\tilde{N}_\mathrm{R})^2}{2m_\mathrm{P}^6} \approx \gamma^2\frac{\tilde{N}_\mathrm{R}^2}{2m_\mathrm{P}^2}, \ (12)$$

$$\eta := m_\mathrm{P}^2\left(\frac{V''}{V}\right) \simeq \gamma + \frac{3\delta\tilde{N}_\mathrm{R}^2}{m_\mathrm{P}^2} \approx \gamma, \quad (13)$$

$$\xi := m_\mathrm{P}^4\left(\frac{V'V'''}{V^2}\right) \simeq \frac{6\delta\tilde{N}_\mathrm{R}^2(\gamma m_\mathrm{P}^2 + \delta\tilde{N}_\mathrm{R}^2)}{m_\mathrm{P}^4}, \quad (14)$$

where prime denotes derivative with respect to N_R. Thus, assuming that the slow-roll approximation is justified (i.e. $\epsilon \ll 1$, $\eta \ll 1$), the spectral index n_s, the tensor-to-scalar ratio $r = A_\mathrm{t}/A_\mathrm{s}$ and the running spectral index $dn_\mathrm{s}/d\ln k$ are given by

$$n_\mathrm{s} \simeq 1 - 6\epsilon + 2\eta \approx 1 + 2\gamma, \quad (15)$$

$$r \simeq 16\epsilon \approx \gamma^2\frac{8\tilde{N}_{\mathrm{R}e}^2}{m_\mathrm{P}^2}, \quad (16)$$

$$\frac{dn_\mathrm{s}}{d\ln k} \simeq 16\epsilon\eta - 24\epsilon^2 - 2\xi \approx -\gamma\frac{12\delta\tilde{N}_{\mathrm{R}e}^2}{m_\mathrm{P}^2}. \quad (17)$$

In the above formulae, $\tilde{N}_{\mathrm{R}e}$ is the field value of the inflaton $\tilde{N}_{\mathrm{R}}$ at $N = 50$ to 70 e-folds before the end of inflation, given approximately by

$$\tilde{N}_{\mathrm{R}e} \approx \tilde{N}_{\mathrm{R}c}\, e^{\gamma N}\,, \tag{18}$$

with the critical value $\tilde{N}_{\mathrm{R}c}$ at the end of inflation defined in Eq. (10). The experimental data on the spectral index from WMAP $n_{\mathrm{s}} = 0.99 \pm 0.04$ [21] restricts γ to be roughly $|\gamma| \lesssim 0.02$. As discussed above, γ controls the sneutrino mass during inflation. In this model it stems mainly from supergravity corrections. Realizing a very small value of γ would require some tuning. In addition, we see that the tensor-to-scalar ratio $r = A_{\mathrm{t}}/A_{\mathrm{s}}$ and the running spectral index $\mathrm{d}n_{\mathrm{s}}/\mathrm{d}\ln k$ are suppressed by higher powers of γ or by $\tilde{N}_{\mathrm{R}}^2/m_{\mathrm{P}}^2$ and are thus generically small. Especially the prediction for the tensor-to-scalar ratio $r \ll \gamma^2$ is thus in sharp contrast to the prediction of $r = 0.16$ for the case of chaotic sneutrino inflation.

In our model, the amplitude of the primordial spectrum is given by

$$P_{\mathcal{R}}^{1/2} \simeq \frac{1}{\sqrt{2\varepsilon}}\left(\frac{H}{2\pi m_{\mathrm{P}}}\right) \approx \frac{\kappa}{2\sqrt{3}\,\gamma\,\pi}\frac{M^2}{m_{\mathrm{P}}\,\tilde{N}_{\mathrm{R}e}}\,. \tag{19}$$

Given the COBE normalization $P_{\mathcal{R}}^{1/2} \approx 5 \times 10^{-5}$ [22], from Eqs. (10), (19) and (18) we obtain

$$\frac{M^2}{M_*\, m_{\mathrm{P}}} \approx 3 \times 10^{-8}\,\frac{\gamma^2\,\sqrt{\beta}}{\kappa\,\lambda_N}\,, \tag{20}$$

which relates the scale M in the superpotential to the cutoff scale M_*. It has to be combined with the constraint $\tilde{N}_{\mathrm{R}e} \ll m_{\mathrm{P}}$ (see Eqs. (18) and (10)) and with $M < M', M_*$.

Reheating and Non-thermal Leptogenesis — In our scenario, the observed baryon asymmetry can arise via non-thermal leptogenesis [8, 9].

Let us assume the situation that the inflaton is the lightest singlet sneutrino $\tilde{N}_1$ and that it dominates leptogenesis and reheating after inflation [23]. This is e.g. the case if the waterfall field ϕ decays earlier than the singlet sneutrino inflaton via heavier singlet neutrinos N_2 (or N_3) with comparably large couplings to ϕ. From Eq. (1), using $\langle\phi\rangle = \sqrt{M'M}$, we see that its mass is given by $M_{\mathrm{R}1} = 2(\lambda_N)_{11}M'M/M_*$ in the basis where the mass matrix M_{R} of the singlet (s)neutrinos is diagonal. It decays mainly via the extended MSSM Yukawa coupling $(Y_\nu)_{i1}\hat{L}_i\hat{H}_{\mathrm{u}}\hat{N}_1$ into slepton and Higgs or into lepton and Higgsino with a decay width given by $\Gamma_{N_1} = M_{\mathrm{R}1}(Y_\nu^\dagger Y_\nu)_{11}/(4\pi)$. The decay of the singlet sneutrino after inflation reheats the universe to a temperature $T_{\mathrm{RH}} \approx (90/(228.75\pi^2))^{1/4}\sqrt{\Gamma_{N_1}m_{\mathrm{P}}}$. If $M_{\mathrm{R}1} \gg T_{\mathrm{RH}}$, the lepton asymmetry is produced via cold decays of the singlet sneutrinos [25]. In this case, the produced baryon asymmetry can be estimated as $n_{\mathrm{B}}/n_\gamma \approx -1.84\,\varepsilon\,T_{\mathrm{RH}}/M_{\mathrm{R}1}$, where ε is the decay asymmetry for the singlet sneutrino decay. Note that ε is bounded by $|\varepsilon| \lesssim \frac{3}{8\pi}\sqrt{\Delta m_{31}^2}M_{\mathrm{R}1}/v_{\mathrm{u}}^2$ [26] for the case of hierarchical singlet (s)neutrinos and light neutrinos. $\Delta m_{31}^2 \approx 2.6 \times 10^{-3}$ eV^2 is the atmospheric neutrino mass squared difference and $v_{\mathrm{u}} = \langle H_{\mathrm{u}}\rangle$. The bound on ε implies $|n_{\mathrm{B}}/n_\gamma| \lesssim 1.84\sqrt{\Delta m_{31}^2}\,T_{\mathrm{RH}}\,3/(8\pi v_{\mathrm{u}}^2)$, and hence $T_{\mathrm{RH}} \gtrsim 10^6$ GeV for the observed baryon-to-photon ratio $n_{\mathrm{B}}/n_\gamma = (6.5^{+0.4}_{-0.8}) \times 10^{-10}$ [21].

To take a concrete example of the above discussion, neutrino Yukawa couplings $(Y_\nu)_{i1} \approx 10^{-6}$ and a sneutrino mass of $M_{\mathrm{R}1} = 10^8$ GeV imply a reheat temperature $T_{\mathrm{RH}} \approx 10^6$ GeV, compatible with the gravitino constraints (see e.g. [10]) in some supergravity theories, saturating the lower bound on T_{RH} [27]. In this example the lightest singlet neutrino is effectively decoupled from the see-saw mechanism as in sequential dominance [29]. Using $M_{\mathrm{R}1} = 2(\lambda_N)_{11}M'M/M_*$ and Eq. (20), taking $\gamma = \beta = 10^{-2}$ yields $\kappa = 10^{-2}M'/M$. In addition, $\tilde{N}_{\mathrm{R}e} \ll m_{\mathrm{P}}$ is satisfied for $MM'/m_{\mathrm{P}}^2 \ll 10^{-4}$ (using Eqs. (18) and (10)). We see that this can be achieved easily with $M \approx 0.1M'$ and both scales somewhat below the GUT scale. M_* is not constrained directly and can for example be around 10^{17} GeV. With $M \approx 10^{15}$ GeV and $M' \approx 10^{16}$ GeV we obtain for instance $\tilde{N}_{\mathrm{R}e} \approx 10^{16}$ GeV, well below the Planck Scale. For $(\lambda_N)_{33} = \mathcal{O}(1)$, the heaviest singlet (s)neutrino has a mass $M_{\mathrm{R}3} \approx 10^{14}$ GeV in this case.

Summary and Conclusions — We have proposed a hybrid inflation scenario in which one (or more) of the singlet sneutrinos, the superpartners of the right-handed neutrinos, play the role of the inflaton. Sneutrinos are present in any extension of the MSSM where the smallness of the observed neutrino masses is explained via the see-saw mechanism. In a minimal model of sneutrino hybrid inflation in supergravity we have shown how the baryon asymmetry of our universe can be explained via non-thermal leptogenesis after inflation. For achieving a low reheat temperature, the Yukawa couplings can have values consistent with first family quark and lepton Yukawa couplings in GUT models. For example, the inflaton can be the lightest singlet sneutrino with a mass around 10^8 GeV and can have Yukawa couplings $\approx 10^{-6}$ in order to realize a reheat temperature $T_{\mathrm{RH}} \approx 10^6$ GeV, consistent with the gravitino constraints in some supergravity theories. In contrast to chaotic inflation, the field values of the singlet sneutrino inflaton in hybrid inflation are well below the Planck scale, so that the supergravity corrections can be carefully monitored. In the minimal model considered here these corrections play an essential role. We have found the spectral index $n_{\mathrm{s}} \approx 1 + 2\gamma$ with $|\gamma| \lesssim 0.02$ and a running spectral index $|\mathrm{d}n_{\mathrm{s}}/\mathrm{d}\ln k| \ll |\gamma|$. Furthermore, sneutrino hybrid inflation predicts a small tensor-to-scalar ratio $r \ll \gamma^2$, much smaller than the prediction $r \approx 0.16$ of chaotic sneutrino inflation. This makes sneutrino hybrid inflation easily distinguishable from chaotic sneutrino inflation.

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[11] We may alternatively consider a superpotential of the form [12]: $\mathcal{W}_{\phi\bar{\phi}} = \kappa \hat{S}(\hat{\phi}^n \hat{\bar{\phi}}^n / M'^{2n-2} - M^2)$, with $n > 1$, where $\phi, \bar{\phi}$ can now transform non-trivially under a continuous symmetry group, e.g. under a global B-L symmetry. The choice $n = 1$ would typically not allow positive square masses for both waterfall fields during inflation. The choice $n > 1$ also allows "smooth" hybrid inflation driven by the S field [13].

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[18] Supergravity corrections to the inflaton mass can easily be of the order of the Hubble parameter, inducing $\eta = O(1)$. In our example, we are going to solve this "η-problem" by explicit tuning. An alternative route could be an effective no-scale supergravity theory leading to zero soft scalar masses at high energy (see e.g. [19]).

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